

BICEP2, the curvature perturbation and supersymmetry

David H. Lyth

Consortium for Fundamental Physics,
Cosmology and Astroparticle Group, Department of Physics,
Lancaster University, Lancaster LA1 4YB, UK

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Abstract

The tensor fraction $r \simeq 0.16$ found by BICEP2 corresponds to a Hubble parameter $H \simeq 1.0 \times 10^{14}$ GeV during inflation. This has two implications for the (single-field) slow-roll inflation hypothesis. First, the inflaton perturbation must account for much more than 10% of the curvature perturbation ζ , which barring fine-tuning means that it accounts for practically all of it. It follows that a curvaton-like mechanism for generating ζ requires an alternative to slow roll such as k-inflation. Second, accepting slow-roll inflation, the excursion of the inflaton field is at least of order Planck scale. As a result, the flatness of the inflaton presumably requires a shift symmetry. I point out that if such is the case, the resulting potential is likely to have at least approximately the quadratic form suggested in 1983 by Linde, which is known to be compatible with the observed r as well as the observed spectral index n_s . The shift symmetry does not require supersymmetry. Also, the big H may rule out a GUT by restoring the symmetry and producing fatal cosmic strings. The absence of a GUT would correspond to the absence of superpartners for the Standard Model particles, which indeed have yet to be found at the LHC. It therefore seems quite possible that the quantum field theory chosen by Nature is not supersymmetric.

Recently, BICEP2 [1] has detected primordial gravitational waves. In this note I discuss two consequences of their result for the generation of the curvature perturbation, and some of their implications.

The BICEP2 measurements gives (after subtracting an estimated foreground) $r = 0.16^{+0.06}_{-0.05}$ where $r = \mathcal{P}_{\text{ten}}(k)/\mathcal{P}_{\zeta}(k)$ is the tensor fraction, evaluated on the scales $x_{\text{ls}}/100 \lesssim k^{-1} \lesssim x_{\text{ls}}$ and $x_{\text{ls}} = 14,000$ Mpc is the distance to the last-scattering surface. (I will call these large scales). I will assume that the tensor perturbation is generated during inflation with Einstein gravity; then

$\mathcal{P}_{\text{ten}}(k) = (8/M_{\text{P}}^2)(H(k)/2\pi)^2$ where $H(k) \equiv \dot{a}/a$ is the Hubble parameter at the epoch of horizon exit $k = aH$.

The spectrum $\mathcal{P}_\zeta(k)$ of the curvature perturbation is measured accurately [2] on scales $10 \text{ Mpc} \lesssim k^{-1} \lesssim x_{\text{ls}}$. (I will call these cosmological scales.) It is nearly constant and equal to $\mathcal{P}_\zeta^{1/2}(k_0) = 4.69 \pm 0.02$ at a ‘pivot scale’ $k_0^{-1} = 20 \text{ Mpc}$. With $r = 0.16$ this corresponds to $H = 1.0 \times 10^{14} \text{ GeV}$, where H without an argument is the large-scale value. This corresponds to an energy scale $\rho^{1/4} = 1.5 \times 10^{16} \text{ GeV}$ which is not far below the Planck scale $M_{\text{P}} = 2.4 \times 10^{18} \text{ GeV}$. With my assumption, four-dimensional field theory is valid up to this scale and the string theory from which it may be derived is relevant only at higher scales even closer to the Planck scale.

For $n_s(k) - 1 \equiv d \ln \mathcal{P}_\zeta / d \ln k$ observation requires $n_s(k_0) - 1 = 0.039 \pm 0.005$ assuming $\alpha_s \equiv n'_s(k_0) = 0$. Assuming just constant α_s doesn’t change the result for $n_s - 1$ much, and gives $\alpha_s(k_0) = -0.014^{+0.016}_{-0.017}$ at 95% confidence.¹ These results suggest that $n_s(k) - 1$ is nearly constant on cosmological scales.²

On the usual assumption that H is nearly constant throughout inflation, the measured value of H determines the number of e -folds of inflation after the scale $k = 1/x_{\text{ls}}$ leaves the horizon given an assumed evolution of the scale factor after inflation. Assuming matter domination until reheating at temperature T_{R} , it is

$$N(1/x_{\text{ls}}) = 61 - \frac{1}{3} \ln \frac{10^9 \text{ GeV}}{T_{\text{R}}}. \quad (1)$$

Requiring only successful BBN one could have $T_{\text{R}} \sim 1 \text{ MeV}$ which would give $N = 47$, but a value closer to 61 is far more likely.

Slow-roll inflation

Slow-roll inflation requires

$$\epsilon(k) \ll 1, \quad |\eta(k)| \ll 1, \quad (2)$$

where $\epsilon(k) = M_{\text{P}}^2(V'/V)^2/2$ and $\eta(k) = V''/V$ evaluated at horizon exit, with $V(\phi)$ is the inflaton potential.³ It also gives $V(N) = 3M_{\text{P}}^2 H^2(N)$ to good accuracy, and

$$H^{-1}(N) dH/dN = \epsilon, \quad \epsilon^{-1} d\epsilon/dN = 2\eta - 4\epsilon, \quad (3)$$

where $N(k)$ is the number of e -folds of inflation after k leaves the horizon. Therefore, $H(N)$ and $\epsilon(N)$ hardly change while $\Delta N = 1$.

¹Except where stated all the other uncertainties are at 68% confidence levels.

²A significant change in $n_s(k)$ has been suggested [1, 3], as one way of resolving the tension between the BICEP2 result and the Planck result [2] $r < 0.11$ (95% confidence). When I come to consider models of slow-roll inflation I will assume that the tension is resolved in some other way.

³I consider only single-field inflation, excluding for example the two-field scenario of [4].

Also, slow-roll gives $n_s(k) - 1 = 2\eta(k) - 6\epsilon(k)$, and $r = 16\epsilon$ where ϵ without an argument is the large-scale value. Finally,

$$d\phi/dN = M_{\text{P}}^2 V'/V = \pm \sqrt{2\epsilon(N)} M_{\text{P}}. \quad (4)$$

Integrating this expression gives

$$N(\phi) = \frac{1}{M_{\text{P}}^2} \int_{\phi_{\text{end}}}^{\phi} \frac{V}{V'} d\phi, \quad (5)$$

where ϕ_{end} is ϕ at the end of inflation. For a non-hybrid model inflation ends when slow-roll fails, ie. when $\max\{\epsilon, |\eta|\} \sim 1$.

The inflaton contribution to ζ

The first implication of $r = 0.16$ for slow-roll inflation concerns the contribution ζ_{inf} of the inflaton perturbation to primordial curvature perturbation ζ [5].⁴ To be more precise, it concerns $\mathcal{P}_{\zeta}^{1/2}(k)$ which is the rms contribution of ζ_{inf} per unit interval of $\ln k$. On large scales slow-roll inflation gives

$$\frac{\mathcal{P}_{\zeta_{\text{inf}}}^{1/2}(k)}{\mathcal{P}_{\zeta}^{1/2}(k)} = \sqrt{\frac{r}{16\epsilon}} \gg \sqrt{\frac{r}{16}} = 0.10. \quad (6)$$

It follows that ζ_{inf} *accounts for much more than 10% of ζ , at least on large scales.*

The requirement $\zeta_{\text{inf}} = \zeta$ was taken to be essential for almost twenty years after the advent of inflation. Then it was realised that all [6] or some [7] of ζ could instead be generated by a curvaton field.⁵ The curvaton acquires its perturbation during inflation but generates \mathcal{P}_{ζ} only at some epoch after inflation when it comes to account for a significant fraction of the energy density. Soon after it was realised that ζ could also be generated by what one might call a modulon field, whose perturbation modulates some process after inflation that would take place anyway even if the modulon didn't exist. The original proposal [9] was to modulate the decay of the inflaton, but many other possibilities have since been considered.

The high value of r found by BICEP2 means that a curvaton or modulon mechanism is practically incompatible with slow-roll inflation, and completely incompatible if the mechanism accounts for practically all of ζ .⁶ That is not a problem for these mechanisms though, because they assume nothing about the mechanism of inflation. Their only input from inflation is the Hubble parameter $H(N)$ while scales of interest are leaving the horizon. Alternatives

⁴Generic r is considered [5], which states a converse equivalent of the present result; that $\mathcal{P}_{\zeta_{\text{inf}}} \ll \mathcal{P}_{\zeta}$ is equivalent to $r \ll 0.1$.

⁵See also [8] for a related scenario involving a bouncing universe.

⁶This is verified in [10] for the linear curvaton model.

to slow-roll are known to exist, such as k-inflation [11], which generate a contribution to ζ that is much smaller than ζ_{inf} allowing a curvaton or modulon scenario.

The inflaton scenario is distinguished from the curvaton and modulon scenarios by the non-gaussianity that it generates. The reduced bi-spectrum $f_{\text{NL}}(k_1, k_2, k_3)$ for the inflaton scenario has a very specific shape [12] which would be a smoking gun for slow-roll inflation, but is only of order 10^{-2} which is probably too small ever to detect. The curvaton and modulon scenarios in contrast generally make f_{NL} constant, and typically of order 1 or bigger which should eventually be detectable. Also, when f_{NL} is constant there is a further prediction $\tau_{\text{NL}} = (6f_{\text{NL}}/5)^2$ for the trispectrum parameter τ_{NL} . If the curvaton or modulon evolves non-linearly f_{NL} need not be constant [13] but is still typically $\gtrsim 1$.

Now we know that $r \simeq 0.16$, *a future detection of $|f_{\text{NL}}(k_1, k_2, k_3)| \gg 10^{-2}$ would strongly suggest that inflation is not slow-roll.* Before, it would just have ruled out $\zeta \simeq \zeta_{\text{inf}}$. Furthermore *an accurate verification of $\tau_{\text{NL}} = (6f_{\text{NL}}/5)^2$ would completely rule out slow-roll inflation.* Before, it would just have required that ζ_{inf} is negligible.

Existing studies of curvaton and modulon scenarios leave $H(N)$ arbitrary. Now that we know H on large scales the predictions are much sharper and should all be revisited. For instance, the axionic curvaton model studied in [14] gives, with the observed H and adopting the simplest version, $f_{\text{NL}} \sim 1$.⁷ The axionic curvaton model is particularly attractive because (i) the flatness of the potential can be protected by a shift symmetry as discussed below for the inflaton potential and (ii) it can generate the observed spectral index even if $\epsilon_H \equiv |dH/dN|/H$ is very small.⁸

The big change in the inflaton field

The second implication of $r \simeq 0.16$ for slow-roll inflation [15, 16] is a lower bound on the change in the inflaton field ϕ during inflation.⁹ It comes from the relation Eq. (4). While large scales leave the horizon the change in N is only $\Delta N \simeq 4$. During this era $\epsilon \geq r/16$, and the corresponding change in ϕ is $\Delta_4\phi \simeq 4\sqrt{2\epsilon}$. It follows that $\Delta\phi$, the total change in ϕ after the scale x_{ls} leaves the horizon, satisfies [15]

$$\Delta\phi \gtrsim 4\sqrt{r/8}M_{\text{P}} = 0.56M_{\text{P}}. \quad (7)$$

In [16] it was pointed out that a stronger result holds for the total change in ϕ if $\epsilon(k)$ doesn't decrease during inflation:

$$\Delta\phi \geq N\sqrt{r/8} = 8.4[N(1/x_{\text{ls}})/60]M_{\text{P}}, \quad (8)$$

⁷One sees this from Figure 1 of the paper.

⁸A curvaton or modulon model gives $n_s = 1 - 2\epsilon_H + (V''/3H^2)$ where V is the curvaton potential.

⁹In [15, 16], generic r was considered. In [15] an old definition of r was used, which is 6.9/8 times the now-standard definition.

where $N \equiv N(1/x_{\text{ls}})$ is the total number of e -folds.¹⁰

The potential of slow-roll inflation

I take the inflaton field ϕ to be canonically normalized, and to be described by an effective field theory valid up to some scale $M \lesssim M_{\text{P}}$ (but bigger than the inflationary energy scale). One might identify M with the string scale.

According to a commonly held view, the tree-level potential will contain all terms allowed by the symmetries. Assuming just symmetry under $\phi \rightarrow -\phi$ we will have

$$V = \frac{1}{2}m^2\phi^2 + \sum \lambda_d \phi^d / M^{d-4}, \quad (9)$$

with the sum over even $d > 2$. Also it is commonly assumed that in the absence of a suitably broken symmetry under $\phi \rightarrow \phi + \text{const}$ (shift symmetry) one will have $|m^2| \sim M^2$ and $|\lambda_d| \sim 1$. But Eq. (9) gives

$$\eta(N) = \frac{m^2}{3H^2} + \frac{\sum d(d-1)\lambda_d \phi^d / M^{d-4}}{3H^2 \phi^2}, \quad (10)$$

which barring cancellations requires $m^2 \ll 3H^2$ and

$$d(d-1)|\lambda_d| \ll \frac{3H^2}{M^2} \left(\frac{M}{\phi}\right)^{d-2} \lesssim \frac{3H^2}{M_{\text{P}}^2} \left(\frac{M}{M_{\text{P}}}\right)^{d-4} \lesssim 6 \times 10^{-9}. \quad (11)$$

One may of course dissent from the common view [20], but many authors have taken it on board and have proposed a shift symmetry to ensure the flatness of the potential. Confining ourselves to the large-field case that is relevant here, three of the proposals [21, 22, 23] produce an approximately quadratic potential, $\phi \propto \phi^2$.¹¹ Two more [26, 27] produce a sinusoidal potential, corresponding to what has been called Natural Inflation [28]. With $\phi = 0$ taken to be a minimum, that too can give an approximately quadratic potential. There is also the ‘monodromy’ scheme [29] that typically gives a potential $\propto \phi^p$ with $p < 2$.

A quadratic potential was suggested by Linde in 1983 [32], and it accounts for the measured values of both r and $n_s(k)$. To be precise it gives

$$r = 0.16(50/N(1/x_{\text{ls}})), \quad n(k) - 1 = -0.04(50/N(k)). \quad (12)$$

with $N(k_0) = N(1/x_{\text{ls}}) - 6.5$. These slow-roll predictions are compatible with current observations. They can be altered slightly while maintaining

¹⁰The result Eq. (7) (for generic r) is usually called the Lyth bound. Sometimes (for instance in [17, 18, 19]) the result Eq. (8) is also called the Lyth bound even though it was obtained by two authors [16]). This is done in [17, 19], where potentials with decreasing ϵ are incorrectly said to violate the finding of [15].

¹¹ N -flation [24] can do the same thing, but not for generic initial conditions even assuming equal masses [25]. I thank C. Gordon for pointing this out.

agreement, by using the sinusoidal potential [30], by allowing all terms up to quadratic [31] or by using the ϕ^p potential of [29].

One can also move away from strict slow roll in at least two ways. First, in the schemes of [21, 29] the potential can have an additional oscillating component. That would give $\mathcal{P}_\zeta(k)$ an oscillating component, and also give a possibly observable f_{NL} with a distinctive shape [33]. Second, the sinusoidal potential makes a coupling $\propto \phi F_{\mu\nu} \tilde{F}^{\mu\nu}$ to a gauge field quite likely, which could have a variety of effects [34],

One can also consider potentials that are completely different from ϕ^2 potential, yet give values for r and n that are agreement with observation [35]. It is not clear though, how such potentials can be the result of a broken shift symmetry.

Supersymmetry

The inflationary energy density scale, $\rho^{1/4} = 1.5 \times 10^{16}$ GeV is the same as the GUT scale M_G . That must be a coincidence though, because M_G represents the vev of the GUT Higgs fields and not the height of their potential. The height of their potential will be some coupling $\lambda \ll 1$ times M_G^4 . The energy scale of GUT inflation models, which generate the inflationary energy density from the GUT Higgs fields, is therefore too low [36] to generate the observed r .

The high inflationary energy density is actually dangerous for a GUT, because it breaks supersymmetry. This will generate contributions typically of order $\pm H^2 \simeq \pm (10^{14} \text{ GeV})^2$ to the masses-squared of the GUT Higgs fields, which may be bigger than their true masses-squared of order $-\lambda M_{\text{GUT}}^2$. In that case, if the generated contributions are positive for at least some of the GUT Higgs field, the GUT symmetry may be at least partially restored during inflation which could produce cosmic strings at the end of inflation that could be forbidden by observation. The idea of a supersymmetric GUT is also endangered by the failure so far of the LHC to find supersymmetric partners for the Standard Model particles.

Of course, these considerations do not rule out a GUT. One can suppose instead that the Standard Model is embedded in split supersymmetry, and that inflation generates GUT Higgs mass-squared no bigger than H leading to possibly observable non-gaussianity [4].

Even if the Standard Model has no supersymmetric partners so that there is no GUT, there could still be supersymmetry broken at a high scale. One might therefore consider supersymmetry as a mechanism for keeping the inflationary potential sufficiently flat. For inflation models with $\phi \ll M_P$ supersymmetry is indeed an attractive way of obtaining the shift symmetry, because in such models the terms with $d > 2$ can be suppressed by the factor $(\phi/M_P)^{d-2}$ leaving only the terms $m^2\phi^2/$ and $\lambda_4\phi^4$ to worry about. The suppression of λ_4 can be achieved by taking the inflaton to be a flat direction of supersymmetry. Supergravity, in which supersymmetry will generally be embedded, generically

makes $|m^2| \sim H^2$ (the η problem [37]) but that can be solved by accepting an order 1 percent fine tuning, or [38] by the imposition of an additional symmetry.

Supersymmetry is far less attractive now that we need $\phi \gtrsim M_{\text{P}}$, because all of the λ_d need to be suppressed. One can achieve this by imposing an exact shift symmetry on the Kahler potential which could give the ϕ^2 potential [39, 40, 41] but one may question the validity of doing that [40, 42] and once it is abandoned the situation for the Kahler potential looks no better than that for the potential itself. One could also obtain the ϕ^2 potential from D -term inflation, by imposing an exact shift symmetry on the gauge kinetic function [43] but that too seems to have no justification.

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